

Towards Heisenberg Limit in Magnetometry with Parametric Down Converted Photons

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Recent theoretical and experimental papers have shown how one can achieve Heisenberg limited measurements by using entangled photons. Here we show how the photons in non-collinear down conversion process can be used for improving the sensitivity of magneto-optical rotation by a factor of four which takes us towards the Heisenberg limit. Our results apply to sources with arbitrary pumping. We also present several generalizations of earlier results for the collinear geometry. The sensitivity depends on whether the two-photon or four-photon coincidence detection is used.

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I. INTRODUCTION

Parametric down conversion is a process that is used to produce light possessing strong quantum features. Photon pairs generated by this process show entanglement with respect to different physical attributes such as time of arrival [1] and states of polarization [2]. They are increasingly being utilized for very basic experiments to test the foundation of quantum mechanics and to do quantum information processing [2, 3, 4]. It is also recognized that entangled photon pairs could be useful in many practical applications in precision metrology involving e.g. interferometry [5, 6, 7, 8], imaging [9, 10], lithography [11, 12, 13, 14] and spectroscopy [15]. There is a proposal [16] to use electromagnetic fields in *NOON* states to improve the sensitivity of measurements by a factor of N . Some implementations of this proposal exist [17]. In particular, the use of photon pairs in interferometers allows phases to be measured to the precision in the Heisenberg limit where uncertainty scales as $1/N$ [18] as compared to the shot noise limit where it scales as $1/\sqrt{N}$. This means that for large number of particles, a dramatic improvement in measurement resolution should be possible.

In this paper we present an analysis of how parametric down converted photons could be very useful in getting better spectroscopic information about the medium. We demonstrate how the improvement in magneto-optical rotation (MOR) of light could be realized by employing two different schemes with collinear and non-collinear down conversion geometry in compared to use of coherent light. We calculate the resolution that can be achieved in the MOR's both by use of coherent light and down converted light. We discuss that the Heisenberg limit [19] could be reached in magnetometry by the use down converted light.

II. MOR USING COHERENT LIGHT SOURCE

Consider a single mode coherent light travelling in the z -direction and a linear isotropic medium made

anisotropic by the application of the magnetic field \mathbf{B} in the z -direction. The incident field can be written in the form

$$\mathbf{E}(z, t) = \exp(-i\omega t + ikz)(\hat{x}\varepsilon_x + \hat{y}\varepsilon_y) + c.c. \quad (1)$$

The medium is described by the frequency and magnetic field dependent susceptibilities $\chi_{\pm}(\omega)$. That means horizontally and vertically polarized components of the incident light will rotate on travelling the medium of length l and the field at the exit can be written as

$$\mathbf{E}(l, t) = \exp(-i\omega t + ikl)(\hat{x}\varepsilon_{xl} + \hat{y}\varepsilon_{yl}) + c.c. \quad (2)$$

The rotation of the horizontal and vertical components can be expressed by the relations

$$\begin{pmatrix} \varepsilon_{xl} \\ \varepsilon_{yl} \end{pmatrix} = R \begin{pmatrix} \varepsilon_x \\ \varepsilon_y \end{pmatrix} \quad (3)$$

where

$$R = e^{i\theta_+} e^{i\frac{\theta}{2}} \begin{pmatrix} \cos \frac{\theta}{2} & -\sin \frac{\theta}{2} \\ \sin \frac{\theta}{2} & \cos \frac{\theta}{2} \end{pmatrix}, \quad (4)$$

$$\theta = kl(\chi_+ - \chi_-), \quad (5)$$

$$\theta_+ = kl\chi_+. \quad (6)$$

The corresponding quantum mechanical description can be obtained by replacing the classical amplitudes ε_x and ε_y by the annihilation operators a_x and a_y respectively. For measurements with coherent sources one can look at the intensities of the x and y components of the output when the input is x polarized with coherent state amplitude α_x (See Fig. 1(a)). Then the measured quantities will be

$$I_{xl} = \langle a_{xl}^\dagger a_{xl} \rangle = |\alpha_x|^2 \cos^2 \frac{\theta}{2}, \quad (7)$$

$$I_{yl} = \langle a_{yl}^\dagger a_{yl} \rangle = |\alpha_x|^2 \sin^2 \frac{\theta}{2}. \quad (8)$$

One can estimate the minimum detectable rotation angle θ_m by looking at the fluctuations ΔN_d in the photon number difference between horizontal and vertical photons, where the number difference operator is given as

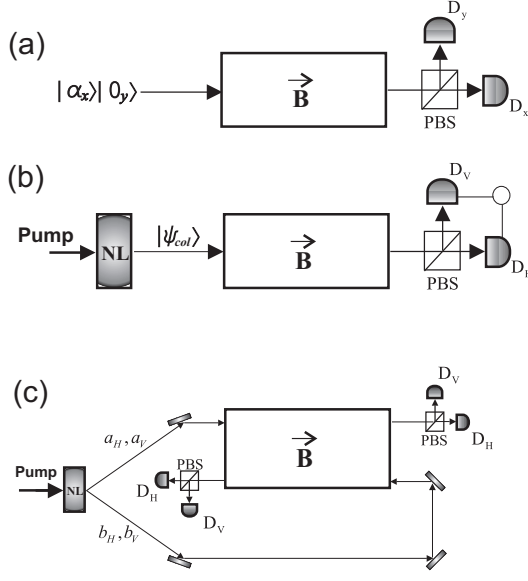


FIG. 1: The setup for the Magneto-optical rotation of light by using (a) coherent source, type-II PDC photons with (b) collinear and (c) non-collinear geometry.

$N_d = a_{yl}^\dagger a_{yl} - a_{xl}^\dagger a_{xl}$. This expression is calculated to be $(\Delta N_d)^2 = |\alpha_x|^2 \sin^2 \theta$ and since the fluctuation noise is 1 we obtain $\theta_m \approx 1/\sqrt{\langle N \rangle}$ where $\langle N \rangle$ is the mean number of input photons which is equal to $|\alpha_x|^2$.

III. MOR USING COLLINEAR TYPE-II PDC AND TWO-PHOTON COINCIDENCE

We now discuss how the results (7) and (8) are modified if we work with down-converted photons. We first consider the collinear case shown in Fig. 1(b). The state produced in collinear PDC can be written by

$$|\psi_{col}\rangle = \frac{1}{\cosh r} \sum_{n=0}^{\infty} (-e^{i\phi} \tanh r)^n |n\rangle_H |n\rangle_V. \quad (9)$$

The value of the parameter r and the phase ϕ are related to the pump amplitude of the nonlinear crystal that is used in the down conversion process and the coupling constant between the electromagnetic field and the crystal. Note that the state $|\psi_{col}\rangle$ is a superposition of n photon pairs of horizontally and vertically polarized modes. Inside the medium, these modes rotate with the same rotation matrix R given in Eq. (4):

$$\begin{pmatrix} a_{Hl} \\ a_{Vl} \end{pmatrix} = R \begin{pmatrix} a_H \\ a_V \end{pmatrix}. \quad (10)$$

One can measure the intensity of each mode:

$$\begin{aligned} I_H &\equiv \langle a_{Hl}^\dagger a_{Hl} \rangle = \sinh^2 r \\ &= \langle a_{Vl}^\dagger a_{Vl} \rangle \equiv I_V. \end{aligned} \quad (11)$$

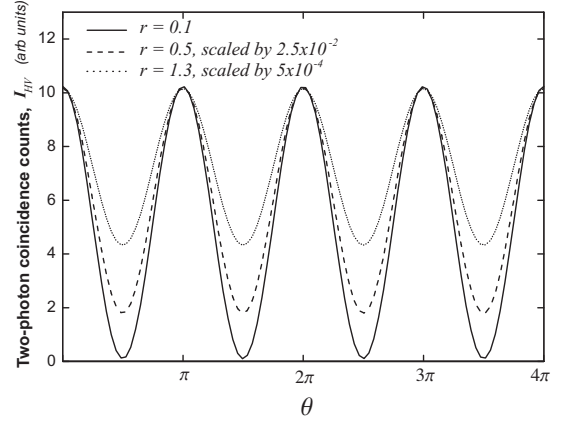


FIG. 2: The MOR plot of two-photon coincidence counts defined by the Eq. (12) in collinear type-II PDC. r is the interaction parameter that defines the pumping strength used in the production of down converted photons and $\theta = kl(\chi_+ - \chi_-)$.

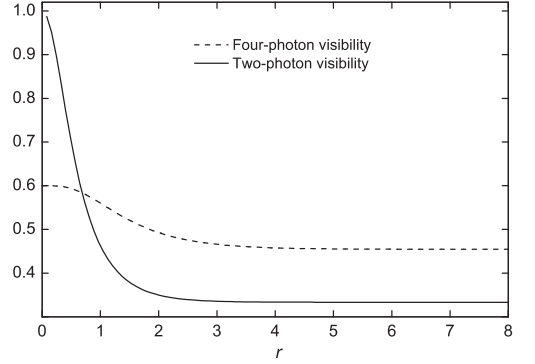


FIG. 3: The visibility of two-photon and four-photon counts defined by the Eqs. (12) and (19) respectively.

And the two-photon coincidence count is:

$$\begin{aligned} I_{HV} &\equiv \langle a_{Hl}^\dagger a_{Vl}^\dagger a_{Hl} a_{Vl} \rangle \\ &= \cos^2 \theta \sinh^2 r \cosh^2 r + \sinh^4 r \end{aligned} \quad (12)$$

Note the difference between Eqs. (7) and (12). With collinearly down-converted photons we measure a rotation angle that is twice as large compared with the angle for a coherent input. For $r \ll 1$ we obtain the same result as given in [15]. The fringe pattern and the visibility is given in Figs. 2 and 3. One can calculate the minimum detectable rotation angle again by looking at the fluctuations in the photon number difference N_d . This is given by $(\Delta N_d)^2 = 4 \sinh^2 r \cosh^2 r \sin^2 \theta = (1 + \langle N \rangle) \langle N \rangle \sin^2 \theta \approx \langle N \rangle^2 \sin^2 \theta$ for large $\langle N \rangle$ where $\langle N \rangle = 2 \sinh^2 r$. Making $(\Delta N_d) \sim 1$ [19] we get $\theta_m \approx 1/\langle N \rangle$. Note that the sensitivity of this quantity is also improved by a factor of $1/\sqrt{\langle N \rangle}$.

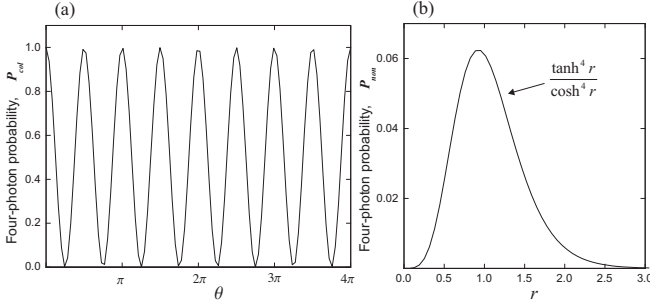


FIG. 4: (a) The normalized four-photon probability defined in Eq. (17), and (b) its envelope wrt the interaction parameter r in the non-collinear geometry.

IV. MOR USING NON-COLLINEAR TYPE-II PDC AND FOUR-PHOTON COINCIDENCE

Next, we discuss the non-collinear PDC case. We have found an arrangement shown in Fig. 1(c) which is especially attractive for improving sensitivity. The entangled photons are coming in two different spatial modes, a and b . While one mode (say a) is going parallel to \mathbf{B} inside the medium, the other is going anti-parallel to it. At the exit we separate the H and V modes by polarizing beam splitters. The state of the input photons can be written in the form [20]

$$|\psi_{non}\rangle = \frac{1}{\cosh^2 r} \sum_{n=0}^{\infty} \sqrt{n+1} (\tanh r)^n |\psi_n\rangle, \quad (13)$$

where

$$|\psi_n\rangle = \frac{1}{\sqrt{n+1}} \sum_{m=0}^n (-1)^m |n-m\rangle_{aH} |m\rangle_{aV} |m\rangle_{bH} |n-m\rangle_{bV}. \quad (14)$$

Here $|m\rangle_{aV}$ represents m vertically polarized photons in mode a . Inside the medium, “+” and “-” polarization components of the modes a and b gain phases $kl\chi_+$ and $kl\chi_-$ respectively. Thus we can write an effective Hamiltonian for the evolution of the state $|\psi_{non}\rangle$ inside the medium as follows:

$$H_{med} = \chi_+ a_+^\dagger a_+ + \chi_- a_-^\dagger a_- - \chi_+ b_+^\dagger b_+ - \chi_- b_-^\dagger b_-, \quad (15)$$

where

$$a_{\pm} = \frac{1}{\sqrt{2}}(a_H \pm ia_V), \quad b_{\pm} = \frac{1}{\sqrt{2}}(b_H \pm ib_V). \quad (16)$$

$$\begin{aligned} I_{HHVV} &= \langle a_H^\dagger a_V^\dagger a_H^\dagger a_V^\dagger \rangle \\ &= (3\cos^2\theta - 1)^2 \sinh^4 r \cosh^4 r + 4(3\cos^2\theta + 1) \sinh^6 r \cosh^2 r + 4\sinh^8 r. \end{aligned} \quad (19)$$

The plot of this quantity for different values of the inter-

action parameter r and the visibility are shown in Figs.

$$\begin{aligned} P_{non} &= |\langle 1_{aH} 1_{aV} 1_{bH} 1_{bV} | \exp(-itH_{med}) |\psi_{non}\rangle|^2 \\ &= \frac{\tanh^4 r}{\cosh^4 r} \cos^2(2\theta), \end{aligned} \quad (17)$$

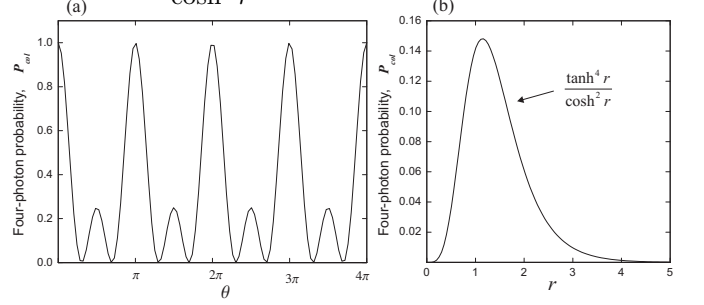


FIG. 5: (a) The normalized four-photon probability defined in Eq. (18), and (b) its envelope wrt the interaction parameter r at the exit ports of PBS in the collinear geometry.

where t is the duration for the state to evolve inside the medium. Note that this four-photon probability has the rotation angle that is four-times as large compared with the angle for a coherent input. The fringe pattern with respect to θ and the probability distribution with respect to r are shown in Fig. 4 (a) and (b).

Next we also examine the four-photon probability in the collinear case. The probability of finding two H -photons and two V -photons at the exit ports of the polarizing beam splitter is given by:

$$\begin{aligned} P_{col} &= |\langle 2_{aH} 2_{aV} | \exp(-itH_{med}) |\psi_{col}\rangle|^2 \\ &= \frac{\tanh^4 r}{\cosh^2 r} \frac{1}{16} [1 + 3\cos(2\theta)]^2, \end{aligned} \quad (18)$$

where we take $H_{med} = \chi_+ a_+^\dagger a_+ + \chi_- a_-^\dagger a_-$ because of the collinear geometry. The normalized plot of this quantity with respect to the magneto-optical rotation angle θ and the envelope of the probability with respect to r are shown in Fig. 5 (a) and (b).

On the other hand one can also calculate the coincidence counts of four photons two-by-two at each detector as given by Glauber's higher order correlation functions:

The plot of this quantity for different values of the inter-

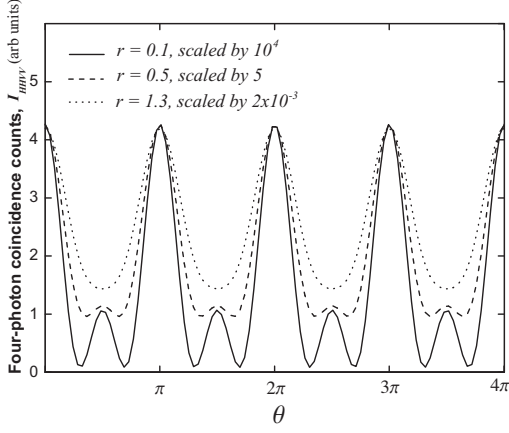


FIG. 6: Four-photon coincidence counts defined in Eq. (19) with different interaction parameter values in the collinear geometry.

6 and 3. Note the distinction between Eqs. (18) and (19) which is a reflection of what the detector is set to measure as we explain now. The former is the probability of the state $|\psi_{col}(t)\rangle$ to be projected onto the particular four-photon subspace $|22\rangle$ i.e. $Tr[|22\rangle\langle 22|\rho_{col}(t)]$ where $\rho_{col}(t) = U|\psi_{col}\rangle\langle\psi_{col}|U^\dagger$ and U is the unitary operator that represents the evolution of the state by the Hamiltonian H_{med} in the collinear geometry. On the other hand, coincidence counting of four-photons at the detectors D_H and D_V (see Fig. 1(b)) is represented by the expectation value $\langle a_H^{\dagger 2} a_V^{\dagger 2} a_H^2 a_V^2 \rangle = Tr[a_H^{\dagger 2} a_V^{\dagger 2} a_H^2 a_V^2 \rho_{col}(t)]$. Note here that the operator $a_H^{\dagger 2} a_V^{\dagger 2} a_H^2 a_V^2$ has the spectral decomposition $\sum_{nm} C_{nm} |nm\rangle\langle nm|$ and obviously it contains the projectors of all $(n+m)$ -photon subspaces with nonzero coefficients C_{nm} . Therefore four-photon counting process at detectors includes not only $|22\rangle$ but all other states $|nm\rangle$ in $|\psi_{col}(t)\rangle$. Here the state $|nm\rangle$ represents n and m photons in the a_H and a_V modes respectively.

V. CONCLUSION

We showed that the use of non-collinear type-II PDC light in MOR's increases the sensitivity by a factor of

four in comparison to coherent light. We also give an argument that minimum rotation uncertainty scales to Heisenberg limit by the use of down converted photons. It should be noted that Heisenberg limit should be understood as an approximate limit at a large mean photon number, that is, the rotation uncertainty approaches the order of $1/\langle N \rangle$ for large $\langle N \rangle$ [8]. The regime with an interaction parameter value of $r = 1.3$ has already been reached in the experiment [21] giving entanglement of 12 photons and an evidence also was given for entanglement up to 100 photons.

APPENDIX A: FOUR-PHOTON PROBABILITY

In this appendix we show the details of the calculation leading to the result given in Eq. (17). One can obtain the result first by solving the Schrödinger equation for the state $|1_{a_H} 1_{a_V} 1_{b_H} 1_{b_V}\rangle$ in the four-photon subspace of the electromagnetic field and having the inner product with the state $|\psi\rangle_{non}$. Since the parts of the Hamiltonian having a and b modes commute, we can solve the Schrödinger equation for the states $|1_{a_H} 1_{a_V}\rangle$ and $|1_{b_H} 1_{b_V}\rangle$ separately. Let us start with a general time-dependent state in the a_H and a_V modes which contain two photons totally;

$$|\phi(t)\rangle = c(t)|20\rangle + d(t)|02\rangle + f(t)|11\rangle \quad (A1)$$

with the initial condition $|\phi(0)\rangle = |11\rangle$. Solving the Schrödinger equation by using the effective Hamiltonian $H = \chi_+ a_+^\dagger a_+ + \chi_- a_-^\dagger a_-$ gives us the result

$$|\phi(t)\rangle = e^{-it\chi} \left[\frac{1}{\sqrt{2}} \sin(\Omega t) |20\rangle - \frac{1}{\sqrt{2}} \sin(\Omega t) |02\rangle + \cos(\Omega t) |11\rangle \right] \quad (A2)$$

where $\chi = \chi_+ + \chi_-$ and $\Omega = \chi_+ - \chi_-$. For a medium of length l , the angle Ωt corresponds to the MOR angle θ which is given in Eq. (5). The solution for the state $|1_{b_H} 1_{b_V}\rangle$ can be obtained just by replacing θ by $-\theta$ because the direction of propagation of the b modes are opposite to that of a modes inside the medium. This is the reason that the part of the effective Hamiltonian for the b_\pm modes takes minus sign in Eq. (15). Consequently we can write the solution of the Schrödinger equation for the state $|1_{a_H} 1_{a_V} 1_{b_H} 1_{b_V}\rangle$ as:

$$\begin{aligned} \exp(-itH_{medium}) |1_{a_H} 1_{a_V} 1_{b_H} 1_{b_V}\rangle &= \exp(-it\chi) \left[\frac{1}{\sqrt{2}} \sin \theta |20\rangle - \frac{1}{\sqrt{2}} \sin \theta |02\rangle + \cos \theta |11\rangle \right] \\ &\otimes \exp(-it\chi) \left[-\frac{1}{\sqrt{2}} \sin \theta |20\rangle + \frac{1}{\sqrt{2}} \sin \theta |02\rangle + \cos \theta |11\rangle \right]. \end{aligned} \quad (A3)$$

Taking the inner product of this with the state $|\psi\rangle_{non}$ and having the absolute square gives us the result given

in Eq. (17).

The result given in Eq. (18) can be obtained by following the same method given above.

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- [1] S. Friberg, C.K. Hong and L. Mandel, Phys. Rev. Lett. **54**, 2011 (1985).
 - [2] P. G. Kwiat, K. Mattle, H. Weinfurter, A. Zeilinger, A. V. Sergienko and Y. Shih, Phys. Rev. Lett. **75**, 4337 (1995).
 - [3] L. Mandel and E. Wolf, *Coherence and Quantum Optics* (Cambridge U. Press, Cambridge, 1995), Chaps. 12 and 22.
 - [4] A. Zeilinger, Rev. Mod. Phys. **71**, 5288 (1999).
 - [5] C. M. Caves, Phys. Rev. D **23**, 1693 (1981).
 - [6] B. Yurke, S. L. McCall and J. R. Klauder Phys. Rev. A **33**, 4033 (1986).
 - [7] J. P. Dowling, Phys. Rev. A **57**, 4736 (1998).
 - [8] M. J. Holland and K. Burnett, Phys. Rev. Lett. **71**, 1355 (1993).
 - [9] T. B. Pittman, Y. H. Shih, D. V. Strekalov and A. V. Sergienko, Phys. Rev. A **52**, R3429 (1995).
 - [10] A. F. Abouraddy, B. E. A. Saleh, A. V. Sergienko and M. C. Teich, Phys. Rev. Lett. **87**, 123602 (2001).
 - [11] A. N. Boto, P. Kok, D. S. Abrams, S. L. Braunstein, C. P. Williams and J. P. Dowling, Phys. Rev. Lett. **85**, 2733 (2000).
 - [12] G. S. Agarwal, R. W. Boyd, E. M. Nagasako and S. J. Bentley, Phys. Rev. Lett. **86**, 1389 (2001).
 - [13] G. Björk, L. L. Sanchez-Soto and J. Söderholm, Phys. Rev. Lett. **86**, 4516, (2001).
 - [14] M. D'Angelo, M. V. Chekhova and Y. Shih, Phys. Rev. Lett. **87**, 013602 (2001).
 - [15] G.S. Agarwal and M.O. Scully, Opt. Lett. **28**, 462 (2003).
 - [16] H. Lee, P. Kok, N.J. Cerf and J.P. Dowling, Phys. Rev. A **65**, 030101(R) (2002).
 - [17] P. Walther et al., Nature (London) **429**, 158 (2004); M.W. Mitchell, J.S. Lundeen and A.M. Steinberg, Nature (London) **429**, 161 (2004).
 - [18] V. Giovannetti, S. Lloyd and L. Maccone, Science **306**, 1330 (2004).
 - [19] Z.Y. Ou, Phys. Rev. A **55**, 2598 (1997).
 - [20] P. Kok and S.L. Braunstein, Phys. Rev. A **61**, 042304 (2000).
 - [21] H. S. Eisenberg, G. Khoury, G. Durkin, C. Simon and D. Bouwmeester, Phys. Rev. Lett. **93**, 193901 (2004).